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# Probing Intrinsic Charm with Semileptonic $B$ Decays

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## Abstract

We discuss semileptonic  $B$  decays of the form  $B \rightarrow J/\Psi \, e \nu \, X$  as possible probes for intrinsic charm. We calculate the leading order perturbative contribution to the process  $B^- \rightarrow J/\Psi \, e^- \bar{\nu}_e$  and find a branching ratio of  $6.8 \times 10^{-10}$ . We estimate the intrinsic charm contribution and find that it could be larger, although its precise value cannot be calculated. Because the kinematically allowed upper limit for these processes is of order  $10^{-7}$ , and because these channels have a very clean experimental signature, we suggest that they should be studied at the  $B$  factories and at hadron colliders.

# 1 Introduction

Interesting possibilities open up if we observe that the wavefunctions of bound states contain Fock states of arbitrarily high particle number. For example the proton wavefunction might be decomposed as

$$|p\rangle = \Psi_{uud}^p |uud\rangle + \Psi_{uudg}^p |uudg\rangle + \Psi_{uudd\bar{d}}^p |uudd\bar{d}\rangle + \Psi_{uuds\bar{s}}^p |uuds\bar{s}\rangle + \Psi_{uudc\bar{c}}^p |uudc\bar{c}\rangle + \dots, \quad (1)$$

and analogously for the  $B^-$  meson:

$$|B^-\rangle = \Psi_{b\bar{u}} |b\bar{u}\rangle + \Psi_{b\bar{u}g} |b\bar{u}g\rangle + \Psi_{b\bar{u}d\bar{d}} |b\bar{u}d\bar{d}\rangle + \Psi_{b\bar{u}s\bar{s}} |b\bar{u}s\bar{s}\rangle + \Psi_{b\bar{u}c\bar{c}} |b\bar{u}c\bar{c}\rangle + \dots. \quad (2)$$

The higher Fock components arise as quantum fluctuations suppressed by  $M^2$ , where  $M$  is the mass of the fluctuation. The  $c\bar{c}$  Fock component of hadrons is referred to as intrinsic charm (IC) [1].

The possibility of detecting signatures for IC in B-meson physics has been recently discussed in Ref. [2]. The authors propose IC as the explanation for a “slow” (low momentum) bump in the inclusive  $B \rightarrow J/\Psi X$  spectrum and the softness of the  $J/\Psi$  spectrum in  $\Upsilon \rightarrow J/\Psi X$  decays, contrary to expectations from the color octet mechanism. IC appearing in higher Fock components of the B-meson wavefunction can manifest itself in two possible ways. It can operate virtually, in a mediation role, and affect decay processes by providing additional channels for the weak interactions. In this way one may enhance some CKM suppressed  $B$  decays, as it has been recently discussed in Ref. [3]. Another instance of this mediation is given by the hypothesis that the IC component of  $B$  mesons is able to explain the “ $\rho\pi$  puzzle” [4]. IC can also manifest itself in processes in which the  $c\bar{c}$  content of the  $B$  meson produces charmed hadrons in the final state, as discussed in Ref. [5] to account for charm production in deep inelastic scattering.

If we consider purely hadronic decays of the  $B$  meson with just one  $c\bar{c}$  pair in the final state we realize, looking for example at Fig. 1, that  $J/\Psi$  X channels can be easily obtained at tree level via  $V_{cb}$ , without recurring to IC. Therefore IC effects, if present, would compete with a formidable background and would be very hard to identify. The same is true if we are looking for modes with two  $D$  mesons in the final state. In the case when there

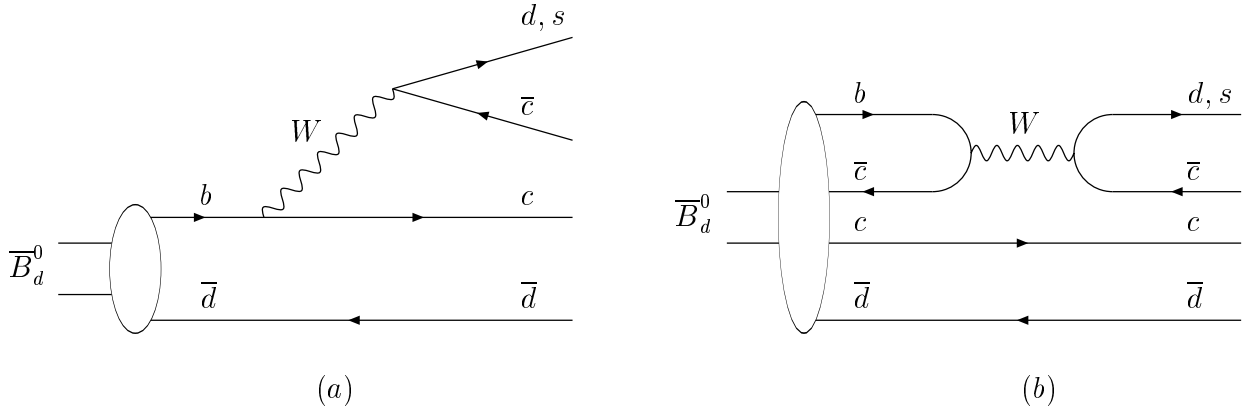


Figure 1: Feynman graphs describing the perturbative contribution to the decay  $\bar{B}_d^0 \rightarrow J/\Psi$  X. (a) is a standard tree-level graph, (b) is an annihilation graph involving IC in the initial state. An IC exchange graph is also allowed.

are three charm (anti)quark particles in the final state (analyzed in Ref. [2]) it is possible to pare down the multiplicity of channels, looking for decays of the type  $B \rightarrow J/\Psi D^{(*)}$  or  $B \rightarrow J/\Psi D^{(*)} X$ . We will return to the former process in our discussion in the final section. In the latter case the phase space available closes rapidly, suppressing the rates of most processes to unobservable levels.

In this paper we discuss semileptonic modes because they present an extremely clean experimental signature. The simplest mode with  $c\bar{c}$  in the final state is  $B^- \rightarrow J/\Psi e^- \bar{\nu}_e$ . Hadron machines like CDF at Fermilab should be able to identify the three outgoing charged

leptons easily. We estimate the size of the perturbative and IC contributions in the following sections.

## 2 The perturbative calculation

In this section we compute the leading contribution to the process  $B^- \rightarrow J/\Psi e^- \bar{\nu}_e$  from two lowest order Feynman diagrams in Fig. 2.

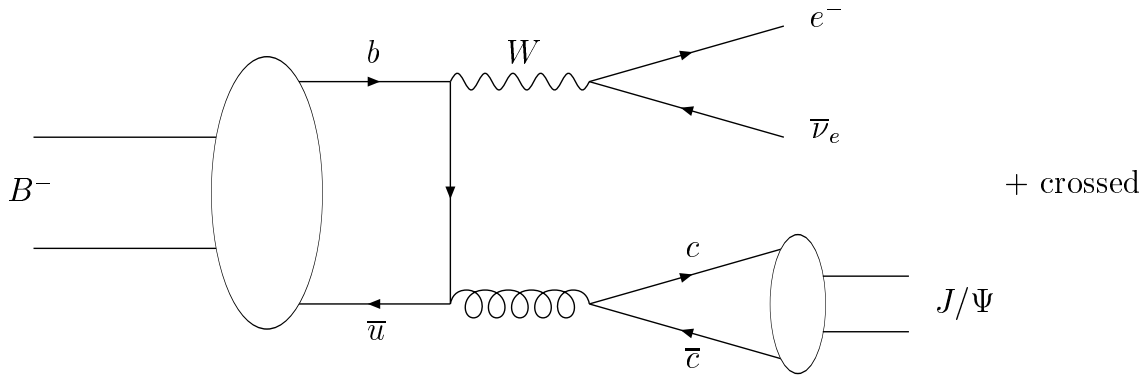


Figure 2: Feynman graphs describing the perturbative contribution to the decay  $B^- \rightarrow J/\Psi e^- \bar{\nu}_e$ .

We adopt the nonrelativistic QCD (NRQCD) factorization formalism [6] to separate  $J/\psi$  formation from production of a  $c\bar{c}$  pair, and express the decay rate in the form

$$\Gamma(B^- \rightarrow J/\Psi e^- \bar{\nu}_e) = \sum_n \Gamma(B^- \rightarrow [c\bar{c}]_n e^- \bar{\nu}_e) \langle \mathcal{O}_n^{J/\Psi} \rangle, \quad (3)$$

where  $\sum_n$  runs over color and angular momentum states of the  $c\bar{c}$  pair and the constants  $\langle \mathcal{O}_n^{J/\Psi} \rangle$  are called NRQCD matrix elements. All the nonperturbative information on the  $J/\Psi$  formation from the  $c\bar{c}$  pair is absorbed into the matrix elements. According to the NRQCD formalism for heavy quarkonium production, the color of the gluon in Fig. 2 is neutralized

by nonperturbative soft gluons in the production process, which are not explicitly shown. For example, there can be soft gluons linking the two blobs in Fig. 2. For the lowest order diagrams in Fig. 2 only a spin-1 and color-octet  $c\bar{c}$  state is able to produce a  $J/\Psi$ , and the corresponding matrix element is  $\langle \mathcal{O}_8^{J/\Psi}({}^3S_1) \rangle$  [7].

Because of the heavy  $J/\Psi$  mass the partonic part of the decay amplitude for  $B^- \rightarrow [c\bar{c}]_8 e^- \bar{\nu}_e$  in Fig. 2 takes place at a short distance  $\sim 1/m_\Psi$ . More precisely, the gluon is off-shell by the  $J/\Psi$  mass, and the denominator of the  $t$ - or  $u$ -channel fermion propagator is at least 5 GeV<sup>2</sup>. Therefore, the partonic part of the decay in Fig. 2 can be factorized from the hadronic matrix element of  $B \rightarrow b\bar{u}$  and calculated perturbatively.

In order to factorize the partonic hard part from the hadronic matrix element of  $B \rightarrow b\bar{u}$ , we can formally write the decay amplitude of  $B \rightarrow g^*W$  in Fig. 2 as:

$$M_i^{\alpha\mu} = \int \frac{d^4 p_b}{(2\pi)^4} \text{Tr} \left[ \hat{T}_B(p_B, p_b) \cdot \hat{H}_i^{\alpha\mu}(p_b, p_\Psi, q) \right], \quad i = 1, 2, \quad (4)$$

where  $\hat{T}_B$  and  $\hat{H}$  represent the matrix element of  $B \rightarrow b\bar{u}$  and partonic part of  $b\bar{u} \rightarrow g^*W$ , respectively, and the partonic momentum  $p_{\bar{u}} = p_B - p_b$ .  $M_1^{\alpha\mu}$  and  $M_2^{\alpha\mu}$  are the contributions from the direct and crossed graphs, respectively.

By decoupling the spinor trace between the matrix element and partonic part, and boosting  $p_B$  to an infinite momentum frame, the decay amplitudes in Eq. (4) can be approximated as

$$M_i^{\alpha\mu} \approx \int dx T_B^\sigma(x) H_{i\sigma}^{\alpha\mu}(p_b = xp_B, p_\Psi, q), \quad (5)$$

where the partonic part  $H$  is given by the lowest order  $b\bar{u} \rightarrow g^*W$  diagrams with  $b$  and  $\bar{u}$

contracted by  $\frac{1}{4}\gamma_5\gamma_\sigma$ , and the matrix element is defined as

$$T_B^\sigma(x) = \int \frac{p_B^+ dy^-}{2\pi} \exp(i x p_B^+ y^-) \langle 0 | \bar{u}(0) \gamma^\sigma \gamma_5 b(y^-) | B^- \rangle , \quad (6)$$

with  $p_B^\mu = (p_B^+, m_B^2/(2p_B^+), 0_\perp)$  using the language of the boost-invariant light-cone wave-function.

Because of the heavy quark mass, the typical virtuality of  $b$  and  $\bar{u}$  in the  $B$  meson should be much smaller than the large momentum scale in  $H$ , and therefore we can approximate the parton momenta  $p_b$  and  $p_{\bar{u}}$  in  $H_{i\sigma}^{\alpha\mu}$  to be

$$p_b = x p_B , \quad p_{\bar{u}} = (1 - x) p_B , \quad \text{with } x = \frac{m_b}{m_B} . \quad (7)$$

Such an approximation corresponds to dropping off power corrections suppressed by the hard momentum scale of the partonic part. By fixing  $x$  in  $H_{i\sigma}^{\alpha\mu}$  we decouple the  $x$  integration in Eq. (5) to find

$$\int dx T_B^\sigma(x) = \langle 0 | \bar{u}(0) \gamma^\sigma \gamma_5 b(0) | B^- \rangle \equiv i f_B p_B^\sigma . \quad (8)$$

Substituting Eq. (8) into Eq. (5) we obtain a factorized expression for the decay amplitude of  $B \rightarrow g^* W$ ,

$$M_i^{\alpha\mu} \approx (i f_B p_B^\sigma) H_{i\sigma}^{\alpha\mu} (p_b = x p_B, p_\Psi, q) \Big|_{x=\frac{m_b}{m_B}} . \quad (9)$$

We can further modify the  $B$  decay constant  $f_B(q^2 = 0) = 0.2 \text{ GeV}$  with a form factor  $1/(1 - q^2/m_B^2)$  to take into account its dependence on  $q^2$  ( $q = q_1 + q_2 = q_e + q_{\bar{\nu}}$ ). This produces the hadronic part of the decay amplitude,

$$M_1^{\alpha\mu} = f_B \frac{g_s g_W}{2\sqrt{2}(1 - q^2/m_B^2)} \frac{1}{(p_\Psi - p_{\bar{u}})^2 - m_{\bar{u}}^2} V_{ub} \left[ -i \varepsilon^{\gamma\delta\alpha\mu} p_{B\gamma} (p_\Psi - p_{\bar{u}})_\delta \right]$$

$$+ p_B^\mu (p_\Psi - p_{\bar{u}})^\alpha + p_B^\alpha (p_\Psi - p_{\bar{u}})^\mu - p_B \cdot (p_\Psi - p_{\bar{u}}) g^{\alpha\mu} \Big] , \quad (10)$$

and for the crossed diagram:

$$\begin{aligned} M_2^{\alpha\mu} = & f_B \frac{g_S g_W}{2\sqrt{2}(1 - q^2/m_B^2)} \frac{1}{(q - p_{\bar{u}})^2 - m_b^2} V_{ub} \Big[ -i \varepsilon^{\gamma\delta\alpha\mu} p_{B\gamma} (p_{\bar{u}} - q)_\delta \\ & + p_B^\mu (-p_{\bar{u}} + q)^\alpha + p_B^\alpha (-p_{\bar{u}} + q)^\mu - p_B \cdot (-p_{\bar{u}} + q) g^{\alpha\mu} \Big] . \end{aligned} \quad (11)$$

In an analogous way, the leptonic part, after taking the modulus squared, yields:

$$L^{\alpha\beta} = g_W^2 \left( \frac{1}{q^2 - m_W^2} \right)^2 \left[ i \varepsilon^{\gamma\delta\alpha\beta} q_{1\gamma} q_{2\delta} + q_1^\beta q_2^\alpha + q_1^\alpha q_2^\beta - q_1 \cdot q_2 g^{\alpha\beta} \right] . \quad (12)$$

The hadronic  $c\bar{c}$  part must contain the projector onto the  $J/\Psi(^3S_1)$  state, plus a color factor which will be taken into account below (see Eq. (16) [7]):

$$J_{\mu\nu}(p_\Psi^2) = \sum_\lambda \left( \frac{-i}{p_\Psi^2} \right) \text{Tr} [(-ig_S \gamma_\mu) \Pi_1^\rho] \varepsilon_\rho^{(\lambda)}(p_\Psi) \cdot \left( \frac{i}{p_\Psi^2} \right) \text{Tr} [(ig_S \gamma_\nu) \Pi_1^{*\sigma}] \varepsilon_\sigma^{*(\lambda)}(p_\Psi) . \quad (13)$$

The projector is defined as [7]

$$\Pi_1^\lambda = \frac{1}{\sqrt{m_\Psi^3}} \frac{1}{4} \left( \not{p}_\Psi - m_\Psi \right) \gamma^\lambda \left( \not{p}_\Psi + m_\Psi \right) . \quad (14)$$

The total contribution from the  $c\bar{c}$  line is then

$$J_{\mu\nu}(p_\Psi^2) = 4 g_S^2 \left( \frac{1}{p_\Psi^2} \right)^2 \frac{1}{m_\Psi} \left[ p_\mu^\Psi p_\nu^\Psi - m_\Psi^2 g_{\mu\nu} \right] . \quad (15)$$

The trace of all the color factors in the graphs for color octet states gives the overall factor [7]

$$\frac{1}{N_c^2 - 1} \sum_{ABC} \text{Tr}[T^B C_8^A] \text{Tr}[T^A C_8^C] = \frac{1}{2} , \quad C_8^A = \sqrt{2} T^A , \quad (16)$$

where  $T^A$ ,  $A = 1\dots 8$ , are the generators of the color group.

As shown in Eq. (3), the squared amplitude for the process  $B^- \rightarrow J/\Psi e^- \bar{\nu}_e$  must include the NRQCD matrix element  $\langle \mathcal{O}_8^{J/\Psi}(^3S_1) \rangle$  defined in [7, 8, 9], which represents the probability for an almost point-like  $c\bar{c}$  pair in a color-octet  $^3S_1$  state to bind and form a  $J/\Psi$ . The matrix element is nonperturbative and has to be extracted from experimental data, and its numerical value is [7, 8]:

$$\langle \mathcal{O}_8^{J/\Psi}(^3S_1) \rangle = 1.06 \times 10^{-2} \text{ GeV}^3 . \quad (17)$$

Putting all this together, the squared amplitude is

$$|\mathcal{M}|^2 = (M_1^{\alpha\mu} + M_2^{\alpha\mu})^* (M_1^{\beta\nu} + M_2^{\beta\nu}) L_{\alpha\beta} J_{\mu\nu}(p_\Psi^2) \frac{1}{2} \langle \mathcal{O}_8^{J/\Psi}(^3S_1) \rangle , \quad (18)$$

including the color factor defined in Eq. (16). Our boson masses are [10]

$$m_B = 5.279 \text{ GeV} , \quad m_\Psi = 3.09687 \text{ GeV} , \quad (19)$$

and we use the constituent masses

$$m_u = 0.35 \text{ GeV} , \quad m_b = 4.9 \text{ GeV} , \quad m_c = m_\Psi/2 , \quad (20)$$

together with

$$|V_{ub}| = 0.003485 [11] , \quad \alpha_S(2m_c) = 0.266 [12] , \quad \tau_{B^-} = 1.653 \times 10^{-12} \text{ sec} [10] . \quad (21)$$

We obtain the following result for the branching ratio:

$$\text{BR}(B^- \rightarrow J/\Psi e^- \bar{\nu}_e) = 6.8 \times 10^{-10} . \quad (22)$$

The differential branching ratio is plotted in Fig. 3.



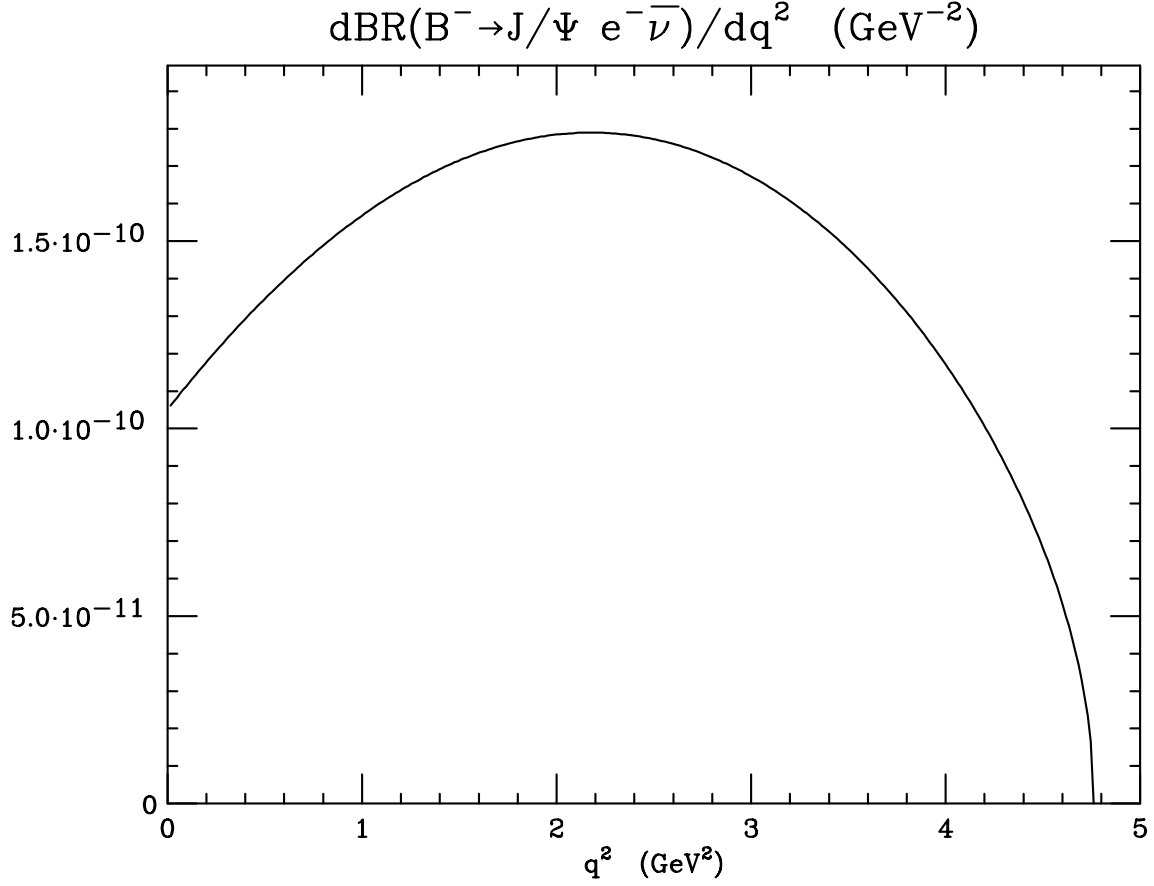


Figure 3: The differential branching ratio for the perturbative contribution to  $B^- \rightarrow J/\Psi e^- \bar{\nu}_e$  is plotted vs.  $q^2 = (q_1 + q_2)^2$ .

To gain some insight into the origin of this result we can approximate the branching ratio assuming a constant  $|\mathcal{M}|^2$ , obtained from estimates of its components discussed above and dimensional analysis. Since

$$d\Gamma(B^- \rightarrow J/\Psi e^- \bar{\nu}_e) = \frac{1}{(2\pi)^3} \frac{1}{32 m_B^3} |\mathcal{M}|^2 dq^2 d(p_\Psi + q_2)^2, \quad (23)$$

and the kinematically allowed area in the phase-space plane  $[q^2, (p_\Psi + q_2)^2]$  is  $A \sim 57 \text{ GeV}^4$ ,

we can write

$$\begin{aligned}
\text{BR}(B^- \rightarrow J/\Psi e^- \bar{\nu}_e) &= \Gamma(B^- \rightarrow J/\Psi e^- \bar{\nu}_e) \frac{\tau_{B^-}}{\hbar} \\
&\approx \frac{1}{(2\pi)^3} \frac{1}{32 m_B^3} A G_F^2 |V_{ub}|^2 (4\pi\alpha_S)^2 \langle \mathcal{O}_8^{J/\Psi}({}^3S_1) \rangle f_B^2 \frac{2}{m_\Psi} \frac{\tau_{B^-}}{\hbar} \\
&\approx 6.2 \times 10^{-10} .
\end{aligned} \tag{24}$$

### 3 The intrinsic charm contribution

The mechanism by which IC in the  $B^-$  meson leads to the decay  $B^- \rightarrow J/\Psi e^- \bar{\nu}_e$  is different. As shown in Fig. 4, the two charm quarks are now part of the initial wavefunction. We picture a  $B^-$  meson in which there is a fluctuation that produces a nearly on-shell  $c\bar{c}$

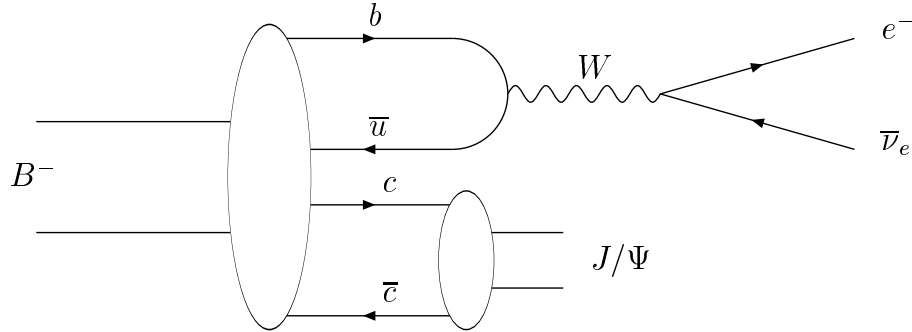


Figure 4: Feynman graph describing the intrinsic charm contribution to the decay  $B^- \rightarrow J/\Psi e^- \bar{\nu}_e$ .

pair that becomes a  $J/\Psi$ , while an off-shell  $b$  quark undergoes the weak decay. Following Ref. [3], the IC effects in  $B$  mesons may be four times larger than in the proton, where it is estimated that they can occur at the 1% level [5]. To obtain our final state we need a

particular  $b\bar{u}c\bar{c}$  configuration that produces a  $J/\Psi$ , *i.e.*  $(c\bar{c})$  in a  $^1S_3$  color singlet state, and  $(b\bar{u})$  in a spin 1 state as well to elude helicity suppression when it decays into the lepton pair  $e\bar{\nu}_e$ . An estimate of the probability of having this desired configuration is model-dependent, and tied to the dynamics of IC. In the most optimistic scenario the fluctuation produces the configuration we need and there is no additional suppression. Alternatively one can estimate the suppression by considering the different combinations of spin, flavor and color which can occur in the wavefunction of the  $B^-$ . For example, from the configurations with the lowest energy listed in Table II of Ref. [3], the suppression factor would be 1/10.

The amplitude can be estimated by treating the charm quarks as spectators. The wavefunction is then the product of a  $(c\bar{c})$  state that generates the  $J/\Psi$  and a virtual  $(b\bar{u})$  state with the same quantum numbers of a  $B^{*-}$ , which decays into leptons. The amplitude squared reduces itself to the product of the leptonic part  $L_{\mu\nu}$  given in Eq. (12) multiplied by the hadronic b–u line contribution:

$$|\mathcal{M}|^2 = -\frac{f_{B^*}^2 m_{B^*}^2}{(1 - q^2/m_{B^*}^2)^2} \frac{g_W^2}{16} |V_{ub}|^2 g^{\mu\nu} L_{\mu\nu} R = 4 \frac{f_{B^*}^2 m_{B^*}^2}{(1 - q^2/m_{B^*}^2)^2} q^2 G_F^2 |V_{ub}|^2 R, \quad (25)$$

where we have modified the decay constant and mass of  $B^{*-}$  with the form factor  $1/(1 - q^2/m_{B^*}^2)$ , and we have included the estimate of Ref. [3] for the relative size of IC,  $R$ , given by

$$R = \frac{1}{(2m_c)^2} \times 0.04. \quad (26)$$

We take  $f_{B^*}(q^2 = 0) = f_B(q^2 = 0) = 0.2$  GeV [13], and  $m_{B^*} = 5.325$  GeV [10]. Finally, we obtain the following branching ratio for the IC component:

$$\text{BR}_{\text{IC}}(B^- \rightarrow J/\Psi e^- \bar{\nu}_e) = 8.65 \times 10^{-9}. \quad (27)$$

We plot the differential branching ratio in Fig. 5. Comparing it with Fig. 3 one notices that the shape of the two distributions is different. This will be useful to separate the perturbative from the IC effect once the experimental sensitivity is sufficient.

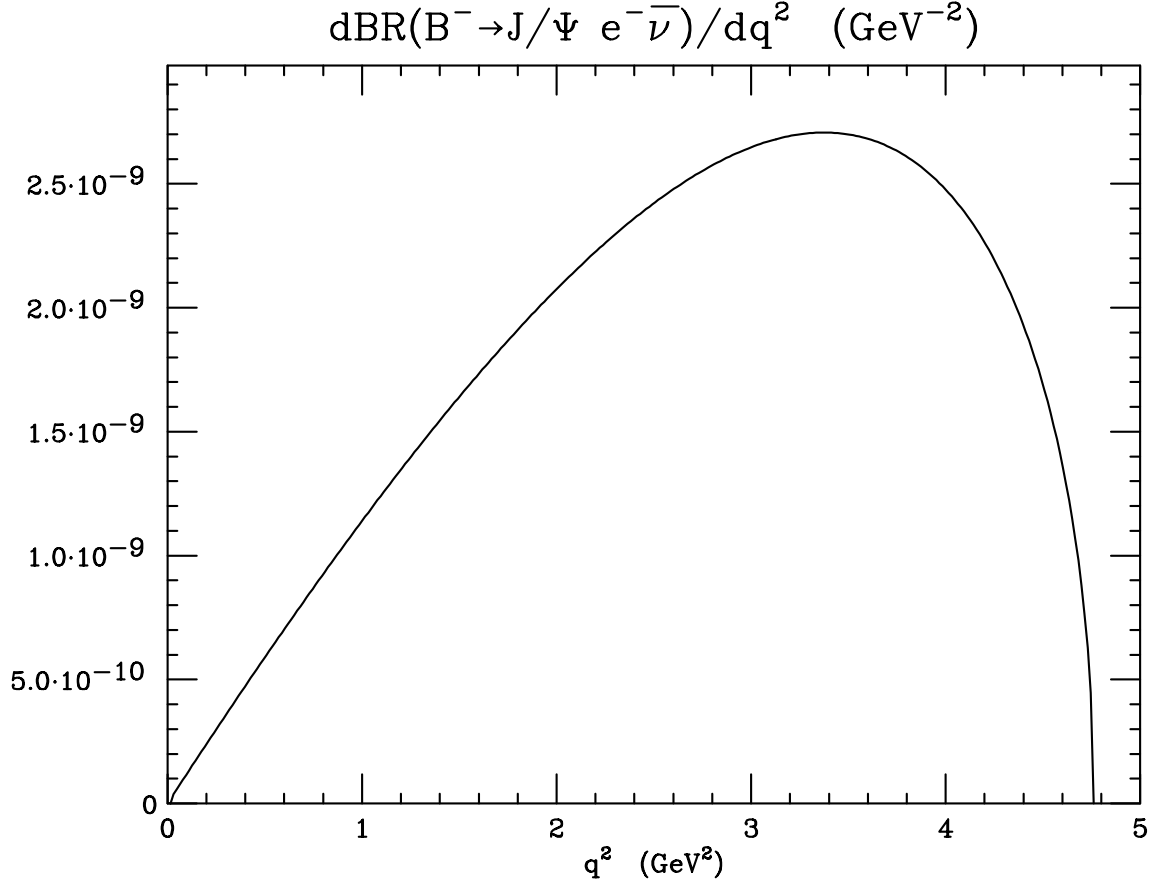


Figure 5: The differential branching ratio for the intrinsic charm contribution to  $B^- \rightarrow J/\Psi e^- \bar{\nu}_e$  is plotted vs.  $q^2 = (q_1 + q_2)^2$ .

The above estimate can also be visualized in terms of an effective vertex as in Fig. 6. The  $J/\Psi$  is directly produced by the fluctuation within the  $B^-$ , and the big dot represents an effective vertex. This vertex should include all the model dependent factors connected with the IC mechanism. An estimate of the maximum size of  $B^- \rightarrow J/\Psi e^- \bar{\nu}_e$  from IC

consists of computing a  $B^* \rightarrow e^- \bar{\nu}_e$  decay with phase space suppression from the additional  $J/\Psi$ . This yields a branching ratio of order  $10^{-7}$ , indicating that it is possible for the IC contribution to be larger than the perturbative result.

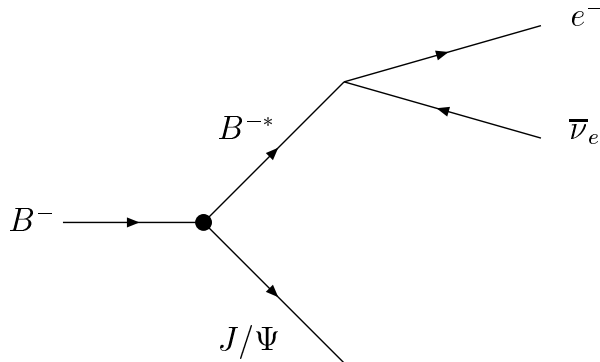


Figure 6: Schematic Feynman graph describing an effective vertex that could model the intrinsic charm contribution to the decay  $B^- \rightarrow J/\Psi e^- \bar{\nu}_e$ .

An alternative estimate for the contribution of IC is obtained by considering the three-body spectator process  $b \rightarrow J/\Psi e^- \bar{\nu}_e X_{u,c}$  depicted in Fig. 7. In this picture we start with a  $B$  wavefunction containing a nearly on-shell (intrinsic)  $c\bar{c}$  pair which directly produces the  $J/\Psi$ . We treat this pair as a spectator. The wavefunction also contains an off-shell  $b$  quark with an effective mass  $m_b^*$  which decays semileptonically:

$$\Gamma_{IC}(b \rightarrow J/\Psi e^- \bar{\nu}_e X_q) \approx \frac{G_F^2 |V_{bq}|^2 m_b^{*5}}{192\pi^3} f\left(\frac{m_q}{m_b^*}\right) \times 0.04, \quad q = u, c, \quad (28)$$

where  $|V_{bu}| = 0.003458$ ,  $|V_{bc}| = 0.041$ . The factor 0.04 is again the estimate of Ref. [3] for the maximum probability of the fluctuation, and

$$f(x) = 1 - 8x^2 + 8x^6 - x^8 - 24x^4 \ln x \quad (29)$$

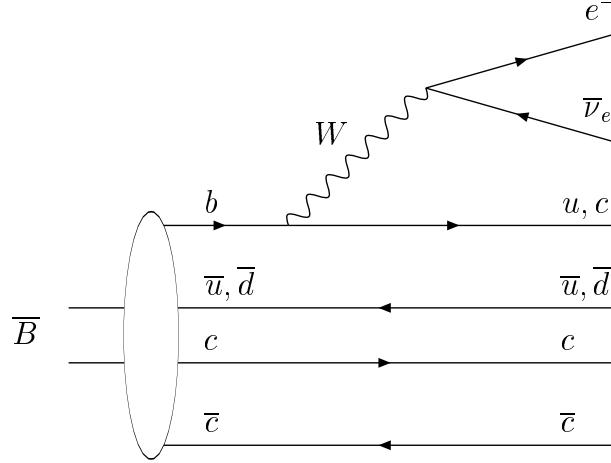


Figure 7: Feynman graphs outlining the IC contribution to the three-body spectator semileptonic decays  $b \rightarrow J/\Psi e^- \bar{\nu}_e X_{u,c}$ .

is the usual phase-space function. The free-quark estimate can be quite misleading for the  $b \rightarrow c$  transition: it indicates that there is enough phase space to produce three charm (anti)quarks in the final state. But the lightest final state with this quark content is  $J/\Psi D$ , so that in fact there will be a large phase space suppression. We illustrate this in Figs. 8 and 9 by plotting the decay rates obtained using the current-quark mass for the c quark,  $m_c = 1.25$  GeV, and a constituent mass  $m_c = 1.6$  GeV, respectively. Because of the strong dependence of the function  $f$  (defined in Eq. (29)) on  $m_c$ , the rate  $\Gamma(b \rightarrow J/\Psi e^- \bar{\nu}_e X_u)$  can be comparable or bigger than  $\Gamma(b \rightarrow J/\Psi e^- \bar{\nu}_e X_c)$  in spite of the suppression due to the unfavorable CKM matrix element  $V_{bu}$ . In the figures  $m_b^*$  varies within the kinematically allowed range  $m_c < m_b^* < m_B - m_\Psi$ . By picking an illustrative value  $m_b^* \sim 2$  GeV we estimate from Fig. 9

$$\text{BR}_{IC}(b \rightarrow J/\Psi e^- \bar{\nu}_e X_q) \approx \begin{cases} 3.7 \times 10^{-7} & \text{for } q = c, \\ 8.9 \times 10^{-7} & \text{for } q = u. \end{cases} \quad (30)$$

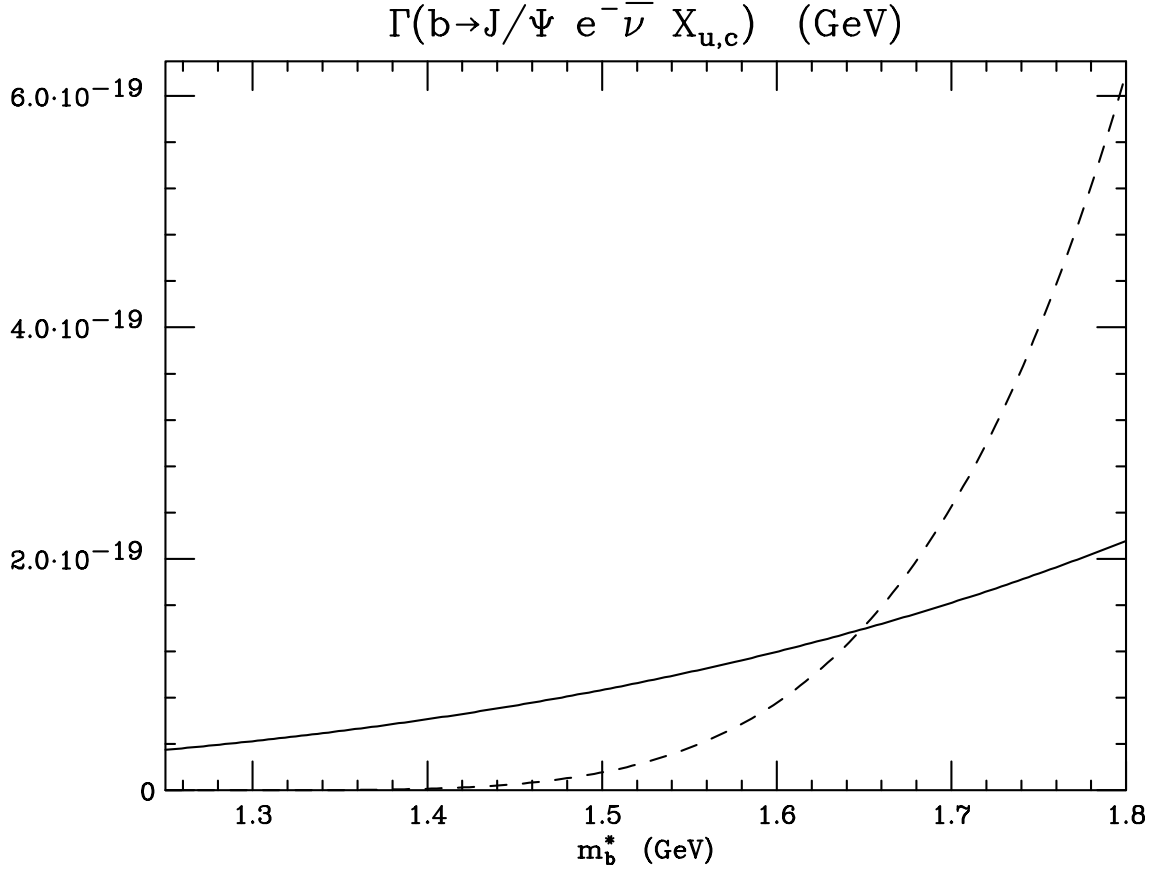


Figure 8: The rates  $\Gamma(b \rightarrow J/\Psi e^- \bar{\nu}_e X_u)$  (solid line) and  $\Gamma(b \rightarrow J/\Psi e^- \bar{\nu}_e X_c)$  (dashed line) are plotted vs. an effective b quark mass  $m_b^*$ , using the current-quark mass for the c quark,  $m_c = 1.25$  GeV, using a free-quark spectator model.

We use a  $b$  quark lifetime  $\tau_b = 1.6 \times 10^{-12}$  sec. With a very clean experimental signature (three charged leptons with two of them reconstructing the  $J/\Psi$  mass) this is a mode well worth looking for.

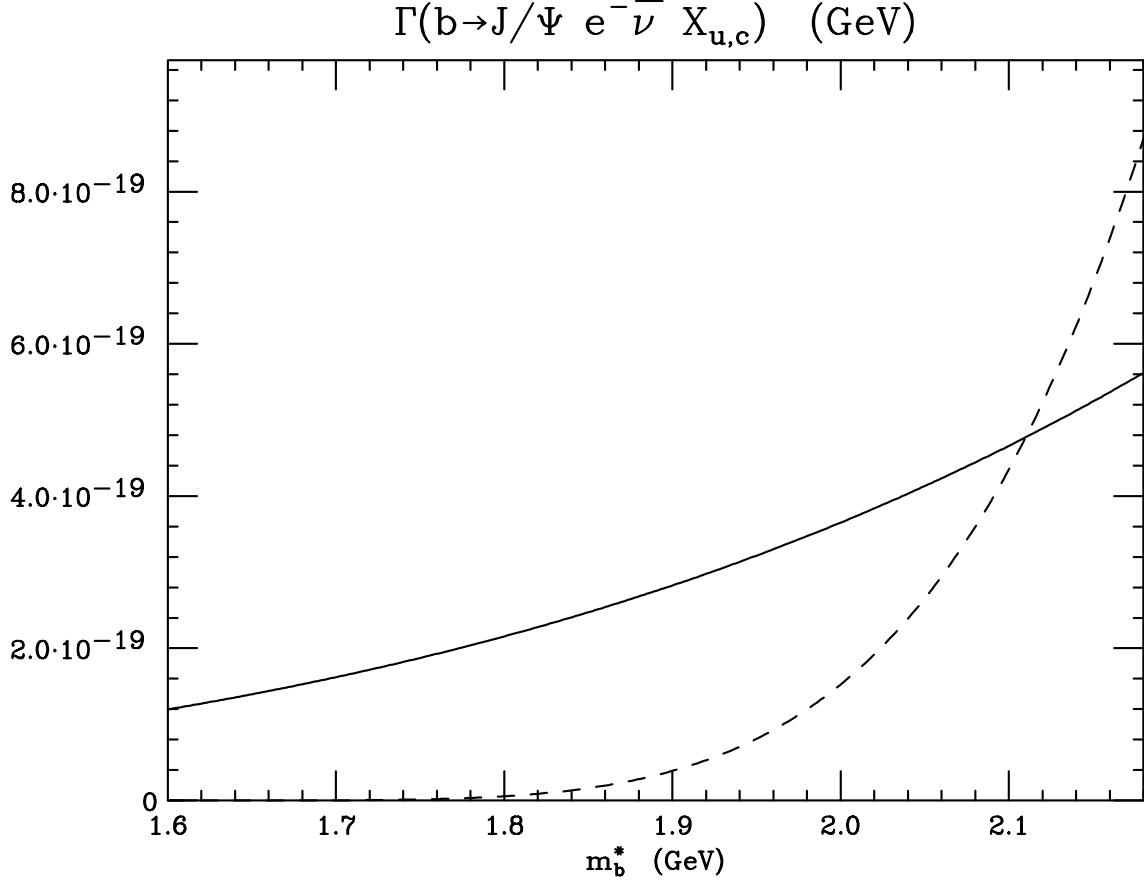


Figure 9: Same as Fig. 8 using a constituent mass for the c quark,  $m_c = 1.6$  GeV.

## 4 Discussion

It is useful to contrast our calculation with that of Ref. [2]. If we accept the current estimates for the branching ratios of  $\overline{B}_d^0 \rightarrow J/\Psi D^{(*)0}$ , *i.e.*  $\text{BR}(\overline{B}_d^0 \rightarrow J/\Psi D^{(*)0}) \sim 10^{-7}$  [2, 12] from PQCD, we are left with the task of evaluating the IC contributions to the above process. This can be done treating the IC pair as spectators (see Fig. 10) and appropriately normalizing the incoming B-meson wavefunction. Analogously to what is shown in Eqs. (1), (2) for the proton and  $B^-$ ,  $\overline{B}_d^0$  has a Fock component  $|\overline{b}d\bar{c}\bar{c}\rangle$  with a coefficient  $\Psi_{\overline{b}d\bar{c}\bar{c}}$ . Again following



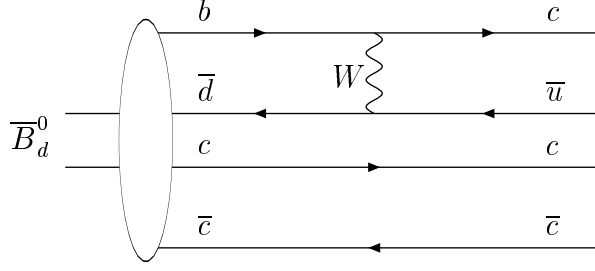


Figure 10: Feynman graph describing the intrinsic charm contribution to the decay  $\overline{B}_d^0 \rightarrow J/\Psi D^{(*)0}$ .

Ref. [3] the probability of such a fluctuation may be as large as 4%. This leads to the following estimate for the branching ratio:

$$\begin{aligned} \text{BR}_{IC}(\overline{B}_d^0 \rightarrow J/\Psi D^{(*)0}) &\approx \frac{1}{N_c^2} \frac{G_F^2 f_B^2 f_D^2}{16\pi m_B^2} \frac{1}{4} (m_B^2 + m_D^2 - m_\Psi^2)^2 |\vec{p}_\Psi| |V_{bc} V_{ud}^*|^2 R \frac{\tau_{B^0}}{\hbar} \\ &= 6.3 \times 10^{-8} . \end{aligned} \quad (31)$$

In the above equation the parameters not yet defined are [10]

$$\begin{aligned} m_{B^0} &= 5.2794 \text{ GeV} , & m_{D^0} &= 1.8645 \text{ GeV} , & f_D &= 0.3 \text{ GeV} , \\ \tau_{B^0} &= 1.548 \times 10^{-12} \text{ sec} , & |V_{ud}| &= 0.975 . \end{aligned} \quad (32)$$

In Eq. (31) we have indicated with  $R$  the expectation of Ref. [3] for the relative size of the IC effect:

$$R = \frac{|\Psi_{b\bar{d}c\bar{c}}|^2}{|\Psi_{b\bar{d}}|^2} = \frac{1}{(2m_c)^2} \times 4\% \approx \frac{1}{m_\Psi^2} \times 0.04 . \quad (33)$$

$N_c = 3$  is the number of colors. If we remove the color suppression factor  $1/N_c^2$  (which arises from a required Fierz rearrangement in the naïve factorization scheme) from our estimate, the result is still comparable to the perturbative one. Our value for the branching ratio in

Eq. (31) is well below the optimistic  $\sim 10^{-5} - 10^{-4}$  expected in Ref. [2] to be experimentally interesting.

At the level of a simple model for specific channels we obtain for  $\text{BR}_{IC}(B^- \rightarrow J/\Psi e^- \bar{\nu}_e) = 8.65 \times 10^{-9}$ , somewhat bigger than the perturbative contribution  $\text{BR}(B^- \rightarrow J/\Psi e^- \bar{\nu}_e) = 6.8 \times 10^{-10}$ , whereas for  $\text{BR}_{IC}(\bar{B}_d^0 \rightarrow J/\Psi D^{(*)0}) \approx 6.3 \times 10^{-8}$  we obtain something comparable to the perturbative result  $\sim 10^{-7}$ .

A comparison between our results and the ones in Ref. [2] may proceed at the level of free-quark estimates in a spectator model. We obtain for  $\text{BR}_{IC}(b \rightarrow J/\Psi e^- \bar{\nu}_e X_q)$  the figures given in Eq. (30), to be compared with the above mentioned result  $10^{-5} - 10^{-4}$ .

In conclusion we have identified a decay channel of the  $B$  meson that has a very clean experimental signature and that could reveal the presence of IC in the  $B$  meson wavefunction. The study of this channel would complement the one suggested in Ref. [2].

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